Are the anti-charmed and bottomed pentaquarks molecular heptaquarks?

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I study the charmed $uudd\bar{c}$ resonance $D^{*-}p$ (3100) very recently discovered by the H1 collaboration at Hera. An anticharmed resonance was already predicted, in a recent publication mostly dedicated to the S=1 resonance $\Theta^+(1540)$. To confirm these recent predictions, I apply the same standard quark model with a quark-antiquark annihilation constrained by chiral symmetry. I find that repulsion excludes the $D^{*-}p$ (3100) as a $uudd\bar{c}$ s-wave pentaquark. I explore the $D^{*-}p$ (3100) as a heptaquark, equivalent to a $N-\pi-D^*$ linear molecule, with positive parity and total isospin $I=0$. I find that the $N-D$ repulsion is cancelled by the attraction existing in the $N-\pi$ and $\pi-D$ channels. In our framework this state is harder to bind than the Θ^+ described by a $k-\pi-N$ borromean bound-state, a lower binding energy is expected in agreement with the H1 observation. Multiquark molecules $N-\pi-D$, $N-\pi-B^*$ and $N-\pi-B$ are also predicted.

I. INTRODUCTION

In this paper I study the anti-charmed $uudd\bar{c}$ resonance $D^{*-}p$ (3100) (narrow hadron resonance of 3099 MeV decaying into a $D^{*-}p$) very recently discovered by the H1 collaboration at HERA [1]. This extends to flavour SU(4) the SU(3) anti-decuplet [2, 3, 4] which includes the recently discovered $\Theta^+(1540)$ [5, 6, 7, 8, 9] and $\Xi^{--}(1860)$ [10, 11, 12]. The $\Theta^+(1540)$, $\Xi^{--}(1860)$ and $D^{*-}p(3100)$ are extremely exciting states, because they may be the first exotic hadrons to be discovered, with quantum numbers that cannot be interpreted as a quark and an anti-quark meson or as a three quark baryon. Exotic multiquarks are expected since the early works of Jaffe [13, 14, 15, 16], and the SU(3) exotic anti-decuplet was first predicted within the chiral soliton model [2, 3, 4]. Pentaquark structures have also been studied in the lattice [17, 18, 19]. However the nature of these particles, their isospin, parity [20] and angular momentum, are yet to be determined.

We recently completed a work on the $uudd\bar{s}$ Θ^+ [21] where we indicate that the Θ^+ is probably a $K-\pi-N$ molecule with binding energy of -30 MeV. In that work we first compute the masses of all the possible s-wave and p-wave $uudd\bar{s}$ pentaquarks, and we verify that these pentaquarks are hundreds of MeV too heavy to explain the Θ^+ resonance, except for the $I=0$, $J^P=1/2^+$ state. However in this channel we find a purely repulsive exotic $N-K$ hard core s-wave interaction. This excludes, in our approach, the Θ^+ as a bare pentaquark $uudd\bar{s}$ state or as a tightly bound s-wave $N-K$ narrow resonance. We then add the $\pi-N$, $\pi-K$ and $N-K$ interactions to study the Θ^+ as a borromean three body s-wave boundstate of a π, a N and a K [21, 25, 26], with positive parity [27] and total isospin $I=0$. In that paper, and in a very recent work [23], we also address the S= -2, Q = -2 state Ξ^{--}, discovered by the NA49 experiment with a mass of 1.862 GeV, indicating that this is a $\bar{K}-N-\bar{K}$ molecule with a binding energy of -60 MeV.

In ref. [21] we also conclude suggesting the existence of similar anti-charmed $uudd\bar{c}$ and anti-bottomed exotic $uudd\bar{b}$ hadrons. The anti-charmed pentaquark was widely expected [22], and its properties will certainly contribute to clarify the nature of the pentaquarks. The state $D^{*-}p$ (3100) may be similar to the Θ^+, with the antiquark $\bar{s}$ replaced by a $\bar{c}$. In this case it is natural to consider replacing the K meson by a D meson or by a D^* meson, because the D^* is also a narrow state. For instance in the new positive parity D_s mesons, [24] the $\bar{K}-D$ and the $\bar{K}-D^*$ multiquarks are respectively candidates to the scalar $D_s(2320)$ and to the axial $D_s(2460)$. The energy of the is $D^{*-}p$ (3100) consistent with a $D^*-\pi-N$ linear molecule with an energy of +15 MeV above threshold. This case differs from the previous ones because here there is no negative binding energy. Nevertheless a system which energy is located slightly above threshold is still a narrow state, and in this sense the $D^{*-}p$ (3100) remains in the same family of the Θ^+ and of the Ξ^{--}.

Moreover a natural theoretical motivation exists for considering heptaquarks (or pentaquarks) and not just p-wave pentaquarks (or baryons) in the exotic multiplet of the Θ^+. Supose that a given s-wave pentaquark hadron H is studied and one concludes that it is unstable. Nevertheless one may consider that a flavor singlet quark-antiquark pair $u\bar{u}+d\bar{d}$ or $s\bar{s}$ is created in the hadron H. When the resulting heptaquark H' remains bound, it is a state with an opposite parity to the original H, where the reversed parity occurs due to the intrinsic parity of fermions and anti-fermions. In this sense the new heptaquark H' can be regarded as the chiral partner of H. And, because H' is expected to be aproximately stable, it is naturally rearranged in a s-wave baryon in two s-wave mesons. The mass of the heptaquark H' is expected to be slightly lower than the exact sum of these standard hadron masses due to the binding energy. This principle explains qualitatively the mass of the Θ^+ [21] for the Ξ^{--} [23] and the masses of the non-exotic multiquarks of the D_s and D_s^* familly [24].

In this paper I extend the techniques used in our first publications to the $D^{*-}p$ (3100) and to the other similar narrow resonances with an anti-charm or anti-bottom

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quark. A standard Quark Model (QM) Hamiltonian is assumed, with a confining potential and a hyperfine term. Moreover the Hamiltonian includes a quark-antiquark annihilation term which is the result of spontaneous chiral symmetry breaking. I start in this paper by reviewing the QM, and the Resonating Group Method (RGM) [28] which is adequate to study states where several quarks overlap. Using the RGM, I show that the corresponding exotic baryon-meson short range s-wave interaction is repulsive in exotic channels and attractive in the channels with quark-antiquark annihilation. The short range repulsion contradicts the existence of narrow pentaquarks with an anti-charm or anti-bottom quark. I proceed with the study of the linear molecules or heptaquarks $D-\pi-N$, $D^*-\pi-N$, $B-\pi-N$ and $B^*-\pi-N$. In particular the total energy of these systems is discussed. Finally I conclude interpreting the $D^{*-}p$ (3100) and predicting related multiquarks.

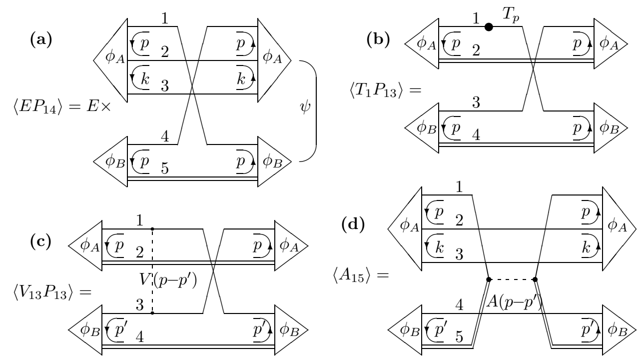

FIG. 1: Examples of RGM overlaps are depicted, in (a) the norm overlap for the meson-baryon interaction, in (b) a kinetic overlap the meson-meson interaction, in (c) an interaction overlap the meson-meson interaction, in (d) the annihilation overlap for the meson-baryon interaction.

II. FRAMEWORK

Our Hamiltonian is the standard QM Hamiltonian,

$$H = \sum_i T_i + \sum_{i<j} V_{ij} + \sum_{i\bar{j}} A_{i\bar{j}} \tag{1}$$

where each quark or antiquark has a kinetic energy T_i with a constituent quark mass, and the colour dependent two-body interaction V_{ij} includes the standard QM confining term and a hyperfine term,

$$V_{ij} = \frac{-3}{16}\vec{\lambda}_i \cdot \vec{\lambda}_j \left[V_{conf}(r) + V_{hyp}(r)\vec{S}_i \cdot \vec{S}_j \right] \ . \tag{2}$$

The QM of eq. (1) reproduces the meson and baryon spectrum with quark and antiquark bound-states (from the heavy quarkonium to the light pion mass). The RGM was first applied by Ribeiro [29] to show that in exotic $N-N$ scattering, the quark-quark potential together with the Pauli repulsion of quarks explains the $N-N$ hard core repulsion. Deus and Ribeiro [30] also showed that, in non-exotic channels, the quark-antiquark annihilation could produce a short core attraction. Recently, addressing a tetraquark system with the $\pi-\pi$ quantum numbers, it was shown that the QM also fully complies with the chiral symmetry, including the Adler zero and the Weinberg theorem [31, 32]. Therefore the QM is adequate to address the anti-decuplet, which was predicted [2, 3, 4] in an effective chiral model.

For the purpose of this paper the details of the potentials in eq. (1) are unimportant, only its matrix elements matter. The hadron spectrum constrains the hyperfine potential,

$$\langle V_{hyp}\rangle \simeq \frac{4}{3}\left(M_\Delta - M_N\right) \simeq M_{K^*} - M_K \ . \tag{3}$$

When a light quark is replaced by a heavy quark, say a charmed quark, the hyperfine interaction is decreased, and it must also be replaced by $\langle V_{hyp\,D}\rangle \simeq M_{D^*} - D_K$. The quark-antiquark annihilation potential $A_{i\bar{j}}$ is also constrained when the quark model produces spontaneous chiral symmetry breaking [33, 34]. The annihilation potential A is present in the π Salpeter equation,

$$\begin{bmatrix} 2T+V & A \\ A & 2T+V \end{bmatrix} \begin{pmatrix} \phi^+ \\ \phi^- \end{pmatrix} = M_\pi \begin{pmatrix} \phi^+ \\ -\phi^- \end{pmatrix} \tag{4}$$

where the π is the only hadron with a large negative energy wave-function, $\phi^- \simeq \phi^+$. In eq. (4) the annihilation potential A cancels most of the kinetic energy and confining potential $2T+V$. This is the reason why the pion has a very small mass. From the hadron spectrum and using eq. (4) the matrix elements of the annihilation potential are determined,

$$\begin{aligned} \langle 2T+V\rangle_{S=0} &\simeq \frac{2}{3}(2M_N - M_\Delta) \\ \Rightarrow \langle A\rangle_{S=0} &\simeq -\frac{2}{3}(2M_N - M_\Delta) \ , \end{aligned} \tag{5}$$

where this result is correct for the annihilation of u or d quarks.

The RGM [28] computes the effective multiquark energy using the matrix elements of the microscopic quark-quark interactions. Any multiquark state can be decomposed in combinations of simpler colour singlets, the baryons and mesons. The wave functions of quarks are arranged in anti-symmetrized overlaps of simple colour singlet hadrons. Once the internal energies E_A and E_B of the two hadronic clusters are accounted,

$$\frac{\langle \phi_b\phi_a| H \sum_p (-1)^p P |\phi_a\phi_b\rangle}{\langle \phi_b\phi_a| \sum_p (-1)^p P |\phi_a\phi_b\rangle} = E_a + E_b + V_{a\,b} \ , \tag{6}$$

where $\sum_p (-1)^p P$ is the anti-symmetrizer, the remaining energy of the meson-baryon or meson-meson system

is computed with the overlap of the inter-cluster microscopic potentials,

$$\begin{aligned} V_{\substack{\text{bar } A \\ \text{mes } B}} &= \langle \phi_B \, \phi_A | - (V_{14} + V_{15} + 2V_{24} + 2V_{25}) 3P_{14} \\ &\quad + 3A_{15} | \phi_A \phi_B \rangle / \langle \phi_B \, \phi_A | 1 - 3P_{14} | \phi_A \phi_B \rangle \\ V_{\substack{\text{mes } A \\ \text{mes } B}} &= \langle \phi_B \, \phi_A | (1 + P_{AB}) [-(V_{13} + V_{23} + V_{14} + V_{24}) \\ &\quad \times P_{13} + A_{23} + A_{14}] | \phi_A \phi_B \rangle \\ &\quad / \langle \phi_B \, \phi_A | (1 + P_{AB})(1 - P_{13}) | \phi_A \phi_B \rangle \ , \end{aligned} \tag{7}$$

where P_{ij} stands for the exchange of particle i with particle j, see Fig. 1. It is clear that quark exchange provides the necessary colour octets to match the Gell-Mann matrices λ_i present in the potential V_{ij}. This results in eq. (3) or eq. (5) times an algebraic colour $\times$ spin $\times$ flavour factor and a geometric momentum overlap [35].

A good approximation for the wave-functions of the ground-state hadrons is the harmonic oscillator wave-function,

$$\phi^{\alpha}_{000}(p_\rho) = \mathcal{N}_\alpha{}^{-1} \exp\left(- \frac{{p_\rho}^2}{2\alpha^2} \right) \ , \ \ \mathcal{N}_\alpha = \left(\frac{\alpha}{2\sqrt{\pi}} \right)^{\frac{3}{2}} \ , \tag{8}$$

where the inverse hadronic radius α can not be estimated by electron-hadron scattering because it is masked by the vector meson dominance. α is the only free parameter in this framework. In the case of vanishing external momenta p_A and p_B, the momentum integral in eq. (7) is simply $\mathcal{N}_\alpha{}^{-2}$.

The annihilation potential only occurs in non-exotic channels. Then it is clear from eq. (5) that the annihilation potential provides an attractive (negative) overlap. The quark-quark(antiquark) potential is dominated by the hyperfine interaction of eq. (3), and in s-wave systems with low spin this results in a repulsive interaction. These results are independent of the details of the quark model that one chooses to consider, provided it is chiral invariant. Therefore I arrive at the attraction/repulsion criterion,

- *whenever the two interacting hadrons have quarks (or antiquarks) with a common flavour, the repulsion is increased by the Pauli principle,*
- *when the two interacting hadrons have a quark and an antiquark with the same flavour, the attraction is enhanced by the quark-antiquark annihilation.*

In the particular case of one nucleon interacting with anti-kaons and with kaons, this implies that the short range exotic $D-N$, D^*-N, $B-N$ and B^*-N interactions are repulsive. This shows that the I=0 s-wave pentaquarks $uudd\bar{c}$ and $uudd\bar{b}$ are certainly quite unstable. Higher isospin or spin systems are certainly more unstable in our framework because they have a higher repulsion. On the other hand the short range $\pi-N$, $\pi-D$, $\pi-D^*$, $\pi-B$ and $\pi-B^*$ interactions can be attractive. This motivates the study of a linear molecule with a N, a π and a D, or a D^+, or a B, or a B^*. Quantitatively [21, 32, 36, 37], the effective potentials computed for the different channels, are

$$\begin{aligned} V_{D-N} &= \frac{1}{2} \frac{\frac{1}{2} + \frac{1}{3}\vec{\tau}_D \cdot \vec{\tau}_N}{\frac{3}{4} - \frac{1}{3}\vec{\tau}_D \cdot \vec{\tau}_N} \langle V_{hyp} \rangle \, \mathcal{N}_\alpha{}^{-2} \\ &\quad + \frac{1}{2} \frac{\frac{1}{2} + \frac{1}{3}\vec{\tau}_D \cdot \vec{\tau}_N}{\frac{3}{4} - \frac{1}{3}\vec{\tau}_D \cdot \vec{\tau}_N} \langle V_{hyp\,D} \rangle \, \mathcal{N}_\alpha{}^{-2} \ , \\ V_{\substack{D-N \\ \to D^*-N}} &= \frac{\frac{1+2\sqrt{3}}{8} + \frac{1+\sqrt{3}}{3}\vec{\tau}_D \cdot \vec{\tau}_N}{\sqrt{3} + \frac{1}{4} + \frac{5}{3}\vec{\tau}_D \cdot \vec{\tau}_N} \langle V_{hyp} \rangle \, \mathcal{N}_\alpha{}^{-2} \\ &\quad + \frac{-\frac{1}{8} - \frac{4}{3}\vec{\tau}_D \cdot \vec{\tau}_N}{\sqrt{3} + \frac{1}{4} + \frac{5}{3}\vec{\tau}_D \cdot \vec{\tau}_N} \langle V_{hyp\,D} \rangle \, \mathcal{N}_\alpha{}^{-2} \ , \\ V_{D^*-N} &= \frac{1}{2} \frac{2 + \frac{7}{3}\vec{\tau}_D \cdot \vec{\tau}_N}{\frac{11}{4} + \frac{7}{3}\vec{\tau}_D \cdot \vec{\tau}_N} \langle V_{hyp} \rangle \, \mathcal{N}_\alpha{}^{-2} \\ &\quad + \frac{1}{2} \frac{\frac{-1}{2} + \frac{5}{3}\vec{\tau}_D \cdot \vec{\tau}_N}{\frac{11}{4} + \frac{7}{3}\vec{\tau}_D \cdot \vec{\tau}_N} \langle V_{hyp\,D} \rangle \, \mathcal{N}_\alpha{}^{-2} \ , \\ V_{\pi-N} &= -\frac{1}{3} \, \vec{\tau}_\pi \cdot \vec{\tau}_N \, \langle A \rangle \, \mathcal{N}_\alpha{}^{-2} \ , \\ V_{\pi-D} &= -\frac{4}{9} \, \vec{\tau}_\pi \cdot \vec{\tau}_D \, \langle A \rangle \, \mathcal{N}_\alpha{}^{-2} \ , \\ V_{\pi-D^*} &= -\frac{4}{9} \, \vec{\tau}_\pi \cdot \vec{\tau}_D \, \langle A \rangle \, \mathcal{N}_\alpha{}^{-2} \ , \end{aligned} \tag{9}$$

where $\vec{\tau}$ are the isospin matrices, normalized with $\vec{\tau}^{\,2} = \tau(\tau+1)$. The vanishing momentum case of eq. (9) is sufficient to compute the scattering lengths with the Born approximation. However the study of binding needs the finite momentum case. In the exotic $D-N$ and D^*-N channels, it can be proved that the geometric result $\mathcal{N}_\alpha{}^{-2}$ is then replaced by the separable interaction $|\phi^\alpha_{000}\rangle\langle\phi^\alpha_{000}|$. In the non-exotic $\bar{\pi}-N$, $\pi-D$ and $\pi-D^*$ channels the present state of the art of the RGM does not allow a precise determination of the finite momentum overlap. Nevertheless I assume for simplicity the same separable interaction. This is a reasonable approximation, because the overlaps decrease when the relative momentum of hadrons A and B increases. Moreover when the hadronic potential V_{AB} is in a separable form $v|\phi_1\rangle\langle\phi_1|$, the energy of the bound state and the matrix element of the potential are simple to compute. A boundstate coincides with a pole in the T matrix at a negative energy,

$$\begin{aligned} T &= |\phi_1 > \frac{v}{1 - g_{0\,11} \, v} < \phi_1| \ , \\ g_{0\,ij} &= < \phi_i | \frac{1}{E + i\epsilon - \frac{p^2}{2\mu}} | \phi_j > \ , \end{aligned} \tag{10}$$

therefore one just has to find the energy that cancels $1 - g_0 \, v$. This method also allows the computation of the matrix element of the potential in the boundstate,

$$< V_{AB} > = v \frac{{g_0}^2}{\sum_j {g_{0\,j1}}^2} \ . \tag{11}$$

When $|\phi_1>$ is the harmonic oscillator state of eq. (8), the necessary condition for binding is,

$$-4\mu \, v \geq \alpha^2 . \tag{12}$$

III. BINDING FLAVOUR $uudd\bar{Q}$ MULTIQUARKS

The simplest pentaquarks are not expected to bind due to the attraction/repulsion criterion. For instance the $D^{*-}p$ (3100) cannot be the ground-state $uudd\bar{c}$ pentaquark because the elementary color singlets $(uud)-(d\bar{c})$ or $(udd)-(u\bar{c})$ are repelled, since the elementary color singlets share the same flavour u or d. This also implies that the D^-p and D^0n systems are unbound. The $uudd\bar{c}$ pentaquarks with spin, flavour or angular momentum excitations will have larger masses and also large widths. Nevertheless the $\bar{c}$ pentaquarks are more subtle than the $\bar{s}$ ones, because the D^* is quite stable when compared with the K^*. Therefore one should also consider to excite the spin in the $l\bar{c}$ cluster, and this amounts to study D^*-N bound-states. Indeed the effective potential of eq. (9) is attractive in this case. However this state is coupled to the $D-N$ case also in eq. (9). Once the coupled channel hamiltonian is diagonalized, the energy of the D^*-N is lifted and the attraction is essentially lost. Therefore the simplest way to have attraction, together with a low energy and with a narrow width, consists in adding at least one quark-antiquark pair to the system.

Then this amounts to include a pion in the system, which can be attracted both by the N baryon and by the D or D^* meson to produce a $N-\pi-D$ or $N-\pi-D^*$ linear molecule. For instance the flavour includes combinations of terms like $uud-d\bar{u}-u\bar{c}$ where the anti-quark $\bar{u}$ in the pion can be annihilated both by the u present in the N and by the u present in the D. According to the attraction/repulsion criterion this produces an attractive interaction. The quark d present in the nucleon cancels only part of the attraction to the pion. Incidently the pion-nucleon $I=1/2$ attraction is fixed by chiral symmetry, see reference [36].

The proposed system $N-\pi-D$ and $N-\pi-D^*$ are similar to the model for $\Theta^+(1540)$ advocated in reference [21], in the present case the anti-quark $\bar{s}$ is replaced by a heavy $\bar{c}$ or $\bar{b}$. The increase of the quark mass does not affect directly the attraction, where the $\bar{Q}$ is just a spectator. However the size of the wave-functions $1/\alpha$ is affected. For instance in an harmonic oscillator potential α is proportional to $\sqrt[4]{\mu}$, and the reduced mass μ doubles when one changes from a light-light meson to a heavy-light meson. This amounts to an increase of nearly 20% of the α in the D or B meson. Because the α parameter is increased only in one of the Jacobi coordinates, the average α in $\pi-D$ or $\pi-D^*$ or $\pi-B$ or $\pi-B^*$ is only expected to suffer a 10%. This increase of α will decrease effectively the attractive interaction. for instance the condition for a 2-body binding in eq. (12) shows that this is equivalent to decrease the strength of the attractive potential by a factor of 20%. Similar results are obtained in different models of confinement, say in the funnel interaction which is more adequate for heavy quarks. In what concerns the repulsive $D-N$ potential, it is decreased in the same way. Moreover the strength of the hyperfine potential is further decreased because $\langle V_{hyp\,D}\rangle << \langle V_{hyp}\rangle$. For example the strength of the repulsive $D-N$ potential is expected to decrease by 30 %.

I now use an adiabatic Hartree method to study the stability of the linear $N-\pi-D$ molecule and related molecules with a D^*, a B or a B^*. Essentially the wave-function of the pion is centered between the nucleon and the D, where the nucleon and the D don't overlap with each other. This results in a linear molecule. for simplicity I use an averaged mass for the nucleon and D and for the pion interaction with these quark clusters. I solve a Schrödinger equation for the nucleon in the potential produced by a pion placed at the origin and by the other heavy-light meson placed at a distance $-\mathbf{a}$ of the pion. The potential of the pion is produced by the D meson at the point $-\mathbf{a}$ and the nucleon $+\mathbf{a}$. This produces three binding energies $E_D\, E_\pi$, E_N, and three wave-functions. In the Hartree method the total energy is the sum of these energies minus the matrix elements of the potential energies, This is easily computed once the two Schrödinger equations are solved, with eqs. (10), (11) and (12). The total energy is a function of the distance $\mathbf{a}$, and I minimize it as a function of a. The same steps are repeated for the $N-\pi-D^*$, $N-\pi-B$ and $N-\pi-B^*$ systems. At the point I am not yet able to bind these linear molecule systems, with a negative binding energy. The same happened for the $N-\pi-K$ when we studied the Θ^+. [21] Nevertheless the picture of a $K-\pi-N$ with a binding energy of 30 MeV is still plausible because the medium range interaction remains to be used. Therefore binding or near-binding is also plausible in the $N-\pi-D$, $N-\pi-D^*$, $N-\pi-B$ and $N-\pi-B^*$ systems.

IV. CONCLUSION AND OUTLOOK

I find that $N-\pi-D$, $N-\pi-D^*$, $N-\pi-B$ and $N-\pi-B^*$ nearly bound I=0 linear s-wave molecules, and positive parity are plausible. The absorption of a low energy pion and the resulting decay into a low energy p-wave $N-D(D^*,B,B^*)$ results in a narrow decay width. The $N-\pi-D^*$ and $N-\pi-B^*$ can also decay respectively into the three body systems $N-\pi-D$ and $N-\pi-D$, but this is again narrow since the $N-D^*$ and $N-B^*$ overlaps are suppressed. Moreover $N-D^*$, and $N-B^*$ pure I=0 s-wave, negative parity, pentaquarks are not completely ruled out, but they couple strongly with the decay channels $N-D$ and $N-B$ (repulsive systems), and this results in a larger width.

Comparing with the $K-\pi-N$ and assuming that it is a bound system, I expect the $N-\pi-D^*$ and $N-\pi-B^*$ to be the most bound systems of the family studied in this paper, with a similar binding energy. Nevertheless the $N-\pi-D^*$ is less bound than the $N-\pi-K$, due to a weaker $\pi-D(D^*,B,B^*)$ attraction when compared with the π_K attraction. Thus an energy of 14 MeV above threshold, corresponding to a $uudd\bar{c}$ mass of 3.099 GeV as observed by the the H1 collaboration is plausible. The

corresponding mass of the multiquark with flavour $uudd\bar{b}$ is of the order of 6.416 GeV. I what concerns the goundstate $N-\pi-D$ and $N-\pi-B$ the binding energy is predicted to be higher because in this case the nucleon-heavy-light meson repulsion is larger. So these states would have an energy some MeV larger than respectively 2.957 GeV and 6.370 Gev.

A more quantitative computation of the masses, sizes, and decay rates of the proposed heptaquarks, including coupled channels and exact three body computations will be done elsewhere. I expect that the most relevant contributions that remain to be included in this framework are the solution of the full three body relativistic Fadeev equations, and the inclusion of the medium range interaction. The medium range interaction, which in nuclear physics is described by the sigma meson exchange, is equivalent to the coupling to channels with multiple pions.

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[1] [H1 Collaboration], arXiv:hep-ex/0403017.
[2] M. Chemtob, Nucl. Phys. **256**, 600 (1985).
[3] M. Praszalowicz, in *Skyrmions and Anomalies*, M. Jezabek and M. Praszalowicz, eds., World Scientific (1987), 112.
[4] D. Diakonov, V. Petrov and M. V. Polyakov, Z. Phys. A **359** 305 (1997) [arXiv:hep-ph/9703373].
[5] T. Nakano *et al.* [LEPS Collaboration], Phys. Rev. Lett. **91**, 012002 (2003) [arXiv:hep-ex/0301020].
[6] V. V. Barmin *et al.* [DIANA Collaboration], Phys. Atom. Nucl. **66**, 1715 (2003) [Yad. Fiz. **66**, 1763 (2003)] [arXiv:hep-ex/0304040].
[7] S. Stepanyan *et al.* [CLAS Collaboration], Phys. Rev. Lett. **91**, 252001 (2003) [arXiv:hep-ex/0307018].
[8] J. Barth *et al.* [SAPHIR Collaboration], arXiv:hep-ex/0307083.
[9] K. T. Knoepfle, M. Zavertyaev and T. Zivko [the HERA-B Collaboration], arXiv:hep-ex/0403020.
[10] C. Alt *et al.* [NA49 Collaboration], arXiv:hep-ex/0310014.
[11] H. G. Fischer and S. Wenig, arXiv:hep-ex/0401014.
[12] J. W. Price, J. Ducote, J. Goetz and B. M. K. Nefkens [CLAS Collaboration], arXiv:nucl-ex/0402006.
[13] R.L. Jaffe, SLAC-PUB-1774, talk presented at the Topical Conf. on Baryon Resonances, Oxford, England, July 5-9, 1976; R. L. Jaffe, Phys. Rev. D **15** 281 (1977).
[14] H. Hogaarsen and P. Sorba, Nucl. Phys. **B145**, 119 (1978).
[15] D. Strottman, Phys. Rev. D **20** 748 (1979).
[16] C. Roiesnel, Phys. Rev. **D20**, 1646 (1979).
[17] F. Csikor, Z. Fodor, S. D. Katz and T. G. Kovacs, JHEP **0311** (2003) 070 [arXiv:hep-lat/0309090].
[18] S. Sasaki, arXiv:hep-lat/0310014.
[19] T. W. Chiu and T. H. Hsieh, arXiv:hep-ph/0403020.
[20] T. Hyodo, A. Hosaka and E. Oset, arXiv:nucl-th/0307105.
[21] P. Bicudo and G. M. Marques, Phys. Rev. D **69**, 011503 (2004) [arXiv:hep-ph/0308073].
[22] M. Karliner and H. Lipkin, hep-ph/0307343; T. Browder, I. Klebanov and D. Marlow, hep-ph/041115; P. Huang, Y. Liu, W. Deng, X. Chen and S. Zhu, hep-ph/0403184; M. Nowak, M. Praszalowicz., M Sadzkowski and J. Wasiluk, hep-ph/0403184; X. He and X. Li, hep-ph/0403191; H. Cheng, C. Chua and C. Huang, hep-ph/0403232; J. Dudek, hep-ph/0403235.
[23] P. Bicudo, arXiv:hep-ph/0403146.
[24] P. Bicudo, arXiv:hep-ph/0401106.
[25] F. J. Llanes-Estrada, E. Oset and V. Mateu, arXiv:nucl-th/0311020.
[26] T. Kishimoto and T. Sato, arXiv:hep-ex/0312003.
[27] Y. s. Oh, H. c. Kim and S. H. Lee, Phys. Rev. D **69**, 014009 (2004) [arXiv:hep-ph/0310019]; Y. s. Oh, H. c. Kim and S. H. Lee, arXiv:hep-ph/0312229.
[28] J. Wheeler, Phys. Rev. **52**, 1083 (1937); **52**, 1107 (1937).
[29] J. E. Ribeiro, Z. Phys. C **5**, 27 (1980).
[30] J. Dias de Deus and J. E. Ribeiro Phys. Rev. D **21**, 1251 (1980).
[31] P. Bicudo, Phys. Rev. C **67**, 035201 (2003).
[32] P. Bicudo, S. Cotanch, F. Llanes-Estrada, P. Maris, J. E. Ribeiro and A. Szczepaniak, Phys. Rev. D **65**, 076008 (2002) [arXiv:hep-ph/0112015].
[33] P. Bicudo and J. E. Ribeiro, Phys. Rev. D **42**, 1611 (1990); 1625 (1990); 1635 (1990).
[34] P. Bicudo, Phys. Rev. C **60**, 035209 (1999).
[35] J. E. Ribeiro, Phys. Rev. D **25**, 2406 (1982); E. van Beveren, Zeit. Phys. C **17**, 135 (1982); J. E. Ribeiro, arXiv:hep-ph/0312035.
[36] P. Bicudo, M. Faria, G. M. Marques and J. E. Ribeiro, to be published in Nuc. Phys. A (2004), arXiv:nucl-th/0106071.
[37] P. Bicudo, arXiv:hep-ph/0401106.